

Scaling Relation for Leptonic Constants of Higher Excitations in Heavy Quarkonium V.V.Kiselev

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Abstract

Introduction

A description of quark bound states demands an application of nonperturbative approaches in QCD, which comes in the strong coupling regime ($\alpha_S \sim 1$) within the region of the hadronization at large distances ($r \sim 1/\Lambda_{QCD}$, $\Lambda_{QCD} \sim 200$ MeV).

Considering the bound states with the heavy quarks ($m_Q \gg \Lambda_{QCD}$), one derives some regularities, simplifying the solution of problem.

For the hadrons with a single heavy quark, the virtualities of the heavy quark are low ($\sim \Lambda_{QCD}$) and one can make the expansion of the heavy quark QCD action over the small parameter Λ_{QCD}/m_Q . In the leading approximation, the effective action possesses the symmetry with respect to the substitution of a heavy quark, moving with a velocity \vec{v} , by any other heavy quark, moving with the same velocity \vec{v} and having an arbitrary orientation of its spin [1]. The symmetry allows one to state the scaling law for leptonic constants of the heavy mesons, containing a single heavy quark, and the universal dependence of the form factors for exclusive semileptonic weak transitions between the heavy hadrons such as $B \rightarrow D^{(*)}l\nu$, so that the universal function has fixed normalization at the zero recoil point.

For the heavy quarkonium ($Q\bar{Q}'$), the nonrelativistic heavy quark motion inside the bound state has allowed one to develop the approach of phenomenological potential models. The QCD-motivated potentials combine the linear rise of the quark interaction energy at the large distances and the Coulomb-like interaction at the small distances [2]. Such models reach the

high accuracy in the description of the mass spectra of heavy quarkonia (the $(c\bar{c})$ charmonium and the $(b\bar{b})$ bottomonium), $\delta m \sim 30$ MeV, however, their accuracy in the description of leptonic constants is very low ($\sim 25\%$).

Another powerful tool for the description of bound states with the heavy quarks has become the QCD sum rules [3], combining perturbative calculations and an account of contributions by the vacuum expectation values of composite operators, i.e. by the quark gluon condensates such as $\langle m\bar{q}q \rangle$, $\langle \alpha_s G_{\mu\nu}^2 \rangle$ and so on. However, making the consideration with the finite number of terms in the QCD perturbation theory for the Wilson's coefficients and taking into the account only the restricted set of quark-gluon condensates, results of the QCD sum rules get unphysical dependence on an external parameter of the sum rule scheme (the number of spectral density moments n or the Borel transformation parameter σ). This dependence essentially decreases the predictive power of QCD sum rules. An additional parameter is also the threshold s_{th} , discriminating the resonant region from the hadronic continuum. Moreover, the weight functions, rapidly dropping with the energy increase, define the averaging scheme of QCD sum rules and do not allow one to draw some conclusions on the contributions by the higher excitations of quarkonium with the given quantum numbers, so that these contributions are practically neglected.

The QCD sum rule scheme, using the data on the quarkonium mass spectrum has been recently offered in ref.[4]. In the scheme, the conditions of nonrelativistic quark motion and the small value of ratio Λ_{QCD}/m_Q have allowed one to state the scaling relation for the leptonic constants of S -wave quarkonia ($Q\bar{Q}$)

$$\frac{f^2}{M} = const. , \quad (1)$$

independently of the heavy quark flavours. The generalization of eq.(1) for the heavy quarkonium ($Q\bar{Q}'$) has been considered in ref.[5], so that

$$\frac{f^2}{M} \left(\frac{M}{4\mu} \right)^2 = const. , \quad (2)$$

where $\mu = m_Q m_{Q'}/(m_Q + m_{Q'})$ is the reduced mass of $(Q\bar{Q}')$ system.

In the present paper we use the QCD sum rule scheme of refs.[4, 5] to de-

rive the relation for the leptonic constants of nS -states of heavy quarkonium

$$\frac{f_{n_1}^2}{f_{n_2}^2} = \frac{n_2}{n_1} , \quad (3)$$

that does not depend on the heavy quarkonium content.

In Section 1 we describe the QCD sum rule scheme, using the spectroscopic data, and derive eq.(3). In Section 2 we make the phenomenological analysis of relation (3) and show, that it gives a good description of the experimental relations for the leptonic constants in the ψ - and Υ -particle families. In the Conclusion we discuss the obtained results.

1 QCD Sum Rules for Heavy Quarkonium

Let us consider the two-point correlator functions of quark currents

$$\Pi_{\mu\nu}(q^2) = i \int d^4x e^{iqx} \langle 0 | T J_\mu(x) J_\nu^\dagger(0) | 0 \rangle , \quad (4)$$

$$\Pi_P(q^2) = i \int d^4x e^{iqx} \langle 0 | T J_5(x) J_5^\dagger(0) | 0 \rangle , \quad (5)$$

where

$$J_\mu(x) = \bar{Q}_1(x) \gamma_\mu Q_2(x) , \quad (6)$$

$$J_5(x) = \bar{Q}_1(x) \gamma_5 Q_2(x) , \quad (7)$$

$$(8)$$

Q_i is the spinor field of the heavy quark with $i = c, b$.

Further, write down

$$\Pi_{\mu\nu} = \left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2} \right) \Pi_V(q^2) + \frac{q_\mu q_\nu}{q^2} \Pi_S(q^2) , \quad (9)$$

where Π_V and Π_S are the vector and scalar correlator functions, respectively. In what follows we will consider the vector and pseudoscalar correlators: $\Pi_V(q^2)$ and $\Pi_P(q^2)$.

Define the leptonic constants f_V and f_P

$$\langle 0 | J_\mu(x) | V(\lambda) \rangle = i \epsilon_\mu^{(\lambda)} f_V M_V e^{ikx} , \quad (10)$$

$$\langle 0 | J_{5\mu}(x) | P \rangle = i k_\mu f_P e^{ikx} , \quad (11)$$

where

$$J_{5\mu}(x) = \bar{Q}_1(x) \gamma_5 \gamma_\mu Q_2(x) , \quad (12)$$

so that

$$\langle 0 | J_5(x) | P \rangle = i \frac{f_P M_P^2}{m_1 + m_2} e^{ikx} , \quad (13)$$

where $|V\rangle$ and $|P\rangle$ are the state vectors of 1^- and 0^- quarkonia, and λ is the vector quarkonium polarization, k is 4-momentum of the meson, $k_{P,V}^2 = M_{P,V}^2$.

Considering the charmonium ($\psi, \psi' \dots$) and bottomonium ($\Upsilon, \Upsilon', \Upsilon'' \dots$), one can easily show that the relation between the width of leptonic decay $V \rightarrow e^+ e^-$ and f_V has the form

$$\Gamma(V \rightarrow e^+ e^-) = \frac{4\pi}{9} e_i^2 \alpha_{em}^2 \frac{f_V^2}{M_V} , \quad (14)$$

where e_i is the electric charge of quark i .

In the region of narrow nonoverlapping resonances, it follows from eqs.(4) - (13) that

$$\frac{1}{\pi} \Im m \Pi_V^{(res)}(q^2) = \sum_n f_{Vn}^2 M_{Vn}^2 \delta(q^2 - M_{Vn}^2) , \quad (15)$$

$$\frac{1}{\pi} \Im m \Pi_P^{(res)}(q^2) = \sum_n f_{Pn}^2 M_{Pn}^4 \frac{1}{(m_1 + m_2)^2} \delta(q^2 - M_{Pn}^2) . \quad (16)$$

Thus, for the observed spectral function one has

$$\frac{1}{\pi} \Im m \Pi_{V,P}^{(had)}(q^2) = \frac{1}{\pi} \Im m \Pi_{V,P}^{(res)}(q^2) + \rho_{V,P}(q^2, \mu_{V,P}^2) , \quad (17)$$

where $\rho(q^2, \mu^2)$ is the continuum contribution, which is not equal to zero at $q^2 > \mu^2$.

Moreover, the operator product expansion gives

$$\Pi^{(QCD)}(q^2) = \Pi^{(pert)}(q^2) + C_G(q^2) \langle \frac{\alpha_S}{\pi} G^2 \rangle + C_i(q^2) \langle m_i \bar{Q}_i Q_i \rangle + \dots , \quad (18)$$

where the perturbative contribution $\Pi^{(pert)}(q^2)$ is labeled, and the nonperturbative one is expressed in the form of sum of quark-gluon condensates with the Wilson's coefficients, which can be calculated in the QCD perturbative theory.

In eq.(18) we have been restricted by the contribution of vacuum expectation values for the operators with dimension $d = 4$. For $C_G^{(P)}(q^2)$ one has, for instance, [3]

$$C_G^{(P)} = \frac{1}{192m_1m_2} \frac{q^2}{\bar{q}^2} \left(\frac{3(3v^2 + 1)(1 - v^2)^2}{2v^5} \ln \frac{1 + v}{1 - v} - \frac{9v^4 + 4v^2 + 3}{v^4} \right), \quad (19)$$

where

$$\bar{q}^2 = q^2 - (m_1 - m_2)^2, \quad v^2 = 1 - \frac{4m_1m_2}{\bar{q}^2}. \quad (20)$$

The analogous formulae for other Wilson's coefficients can be found in Ref.[3]. In what follows it will be clear that the explicit form of coefficients has no significant meaning for the present consideration.

In the leading order of QCD perturbation theory it has been found for the imaginary part of correlator that [3]

$$\Im m\Pi_V^{(pert)}(q^2) = \frac{\tilde{s}}{8\pi s^2} (3\bar{s}s - \bar{s}^2 + 6m_1m_2s - 2m_2^2s) \theta(s - (m_1 + m_2)^2) \quad (21)$$

$$\Im m\Pi_P^{(pert)}(q^2) = \frac{3\tilde{s}}{8\pi s^2} (s - (m_1 - m_2)^2) \theta(s - (m_1 + m_2)^2), \quad (22)$$

where $\bar{s} = s - m_1^2 + m_2^2$, $\tilde{s}^2 = \bar{s}^2 - 4m_2^2s$.

The one-loop contribution into $\Im m\Pi(q^2)$ can be included into the consideration (see, for example, Ref.[3]). However, we note that the more essential correction is that of summing a set over the powers of (α_s/v) , where v is defined in eq.(20) and is a relative quark velocity, and α_s is the QCD interaction constant. In Ref.[3] it has been shown that account of the Coulomb-like gluonic interaction between the quarks leads to the factor

$$F(v) = \frac{4\pi}{3} \frac{\alpha_s}{v} \frac{1}{1 - \exp(-\frac{4\pi\alpha_s}{3v})}, \quad (23)$$

so that the expansion of the $F(v)$ over $\alpha_s/v \ll 1$ restores, precisely, the one-loop $O(\frac{\alpha_s}{v})$ correction

$$F(v) \approx 1 - \frac{2\pi}{3} \frac{\alpha_s}{v} \dots \quad (24)$$

In accordance with the dispersion relation one has the QCD sum rules, which state that, in average, it is true that, at least, at $q^2 < 0$

$$\frac{1}{\pi} \int \frac{\Im m\Pi^{(had)}(s)}{s - q^2} ds = \Pi^{(QCD)}(q^2), \quad (25)$$

where the necessary subtractions are omitted. $\Im m\Pi^{(had)}(q^2)$ and $\Pi^{(QCD)}(q^2)$ are defined by eqs.(15) - (17) and eqs.(18) - (24), respectively. eq.(25) is the base to develop the sum rule approach in the forms of the correlator function moments and of the Borel transform analysis (see Ref.[3]). The truncation of the set in the right hand side of eq.(25) leads to the mentioned unphysical dependence of the $f_{P,V}$ values on the external parameter of the sum rule scheme.

Further, let us use the conditions, simplifying the consideration due to the heavy quarkonium.

1.1 Nonperturbative Contribution

We assume that, in the limit of the very heavy quark mass, the power corrections of nonperturbative contribution are small. From eq.(19) one can see that, for example,

$$C_G^{(P)}(q^2) \approx O\left(\frac{1}{m_1 m_2}\right), \quad \Lambda/m_{1,2} \ll 1, \quad (26)$$

where v is fixed, $q^2 \sim (m_1 + m_2)^2$, when $\Im m\Pi^{(pert)}(q^2) \sim (m_1 + m_2)^2$. It is evident that, due to the purely dimensional consideration, one can believe that the Wilson's coefficients tend to zero as $1/m_{1,2}^2$.

Thus, the limit of very large heavy quark mass implies that one can neglect the quark-gluon condensate contribution.

1.2 Nonrelativistic Quark Motion

The nonrelativistic quark motion implies that, in the resonant region, one has, in accordance with eq.(20),

$$v \rightarrow 0. \quad (27)$$

So, one can easily find that in the leading order

$$\Im m\Pi_P^{(pert)}(s) \approx \Im m\Pi_V^{(pert)}(s) \rightarrow \frac{3v}{8\pi^2} s \left(\frac{4\mu}{M}\right)^2, \quad (28)$$

so that with account of the Coulomb factor

$$F(v) \simeq \frac{4\pi}{3} \frac{\alpha_S}{v}, \quad (29)$$

one obtains

$$\Im m \Pi_{P,V}^{(pert)}(s) \simeq \frac{\alpha_S}{2} s \left(\frac{4\mu}{M} \right)^2 . \quad (30)$$

1.3 "Smooth Average Value" Scheme of the Sum Rules

As for the hadronic part of the correlator, one can write down for the narrow resonance contribution

$$\Pi_V^{(res)}(q^2) = \int \frac{ds}{s - q^2} \sum_n f_{Vn}^2 M_{Vn}^2 \delta(s - M_{Vn}^2) , \quad (31)$$

$$\Pi_P^{(res)}(q^2) = \int \frac{ds}{s - q^2} \sum_n f_{Pn}^2 \frac{M_{Pn}^4}{(m_1 + m_2)^2} \delta(s - M_{Pn}^2) , \quad (32)$$

The integrals in eqs.(31)-(32) are simply calculated, and this procedure is generally used.

In the presented scheme, let us introduce the function of state number $n(s)$, so that

$$n(m_k^2) = k . \quad (33)$$

This definition seems to be reasonable in the resonant region. Then one has, for example, that

$$\frac{1}{\pi} \Im m \Pi_V^{(res)}(s) = s f_{Vn(s)}^2 \frac{d}{ds} \sum_k \theta(s - M_{Vk}^2) . \quad (34)$$

Further, it is evident that

$$\frac{d}{ds} \sum_k \theta(s - M_k^2) = \frac{dn(s)}{ds} \frac{d}{dn} \sum_k \theta(n - k) , \quad (35)$$

and eq.(31) can be rewritten as

$$\Pi_V^{(res)}(q^2) = \int \frac{ds}{s - q^2} s f_{Vn(s)}^2 \frac{dn(s)}{ds} \frac{d}{dn} \sum_k \theta(n - k) . \quad (36)$$

The "smooth average value" scheme means that

$$\Pi_V^{(res)}(q^2) = \langle \frac{d}{dn} \sum_k \theta(n - k) \rangle \int \frac{ds}{s - q^2} s f_{Vn(s)}^2 \frac{dn(s)}{ds} . \quad (37)$$

It is evident that, in average, the first derivative of step-like function in the resonant region is equal to

$$\langle \frac{d}{dn} \sum_k \theta(n - k) \rangle \simeq 1 . \quad (38)$$

Thus, in the scheme one has

$$\langle \Pi_V^{(res)}(q^2) \rangle \approx \int \frac{ds}{s - q^2} s f_{Vn(s)}^2 \frac{dn(s)}{ds} , \quad (39)$$

$$\langle \Pi_P^{(res)}(q^2) \rangle \approx \int \frac{ds}{s - q^2} \frac{s^2 f_{Pn(s)}^2}{(m_1 + m_2)^2} \frac{dn(s)}{ds} . \quad (40)$$

Eqs.(39)-(40) give the average correlators for the vector and pseudoscalar mesons, therefore, due to eq.(25) we state that

$$\Im m \langle \Pi^{(hadr)}(q^2) \rangle = \Im m \Pi^{(QCD)}(q^2) , \quad (41)$$

that gives with account of eqs.(30), (39) and (40) at the physical points $s_n = M_n^2$

$$\frac{f_n^2}{M_n} = \frac{\alpha_S}{\pi} \frac{dM_n}{dn} \left(\frac{4\mu}{M} \right)^2 , \quad (42)$$

where in the limit of heavy quarks we use, that for the resonances one has

$$m_1 + m_2 \approx M , \quad (43)$$

so that

$$f_{Vn} \simeq f_{Pn} = f_n . \quad (44)$$

Thus, one can conclude that for the heavy quarkonia the QCD sum rules give the identity of f_P and f_V values for the pseudoscalar and vector states.

Eq.(42) differs from the ordinary sum rule scheme because it does not contain the parameters, which are external to QCD. The quantity dM_n/dn is purely phenomenological. It defines the average mass difference between the nearest levels with the identical quantum numbers.

Further, as it has been shown in ref.[6], in the region of average distances between the heavy quarks in the charmonium and the bottomonium,

$$0.1 fm < r < 1 fm , \quad (45)$$

the QCD-motivated potentials allow the approximation in the form of logarithmic law [7] with the simple scaling properties, so

$$\frac{dn}{dM_n} = \text{const.} , \quad (46)$$

i.e. the density of heavy quarkonium states with the given quantum numbers do not depend on the heavy quark flavours.

In ref.[5] it has been shown, that relation (46) is also practically valid for the heavy quark potential approximation by the power law (Martin potential) [8], where, neglecting a low value of the binding energy for the quarks inside the quarkonium, one can again get eq.(46).

In ref.[4] it has been found, that relation (46) is valid with the accuracy up to small logarithmic corrections over the reduced mass of quarkonium, if one makes the quantization of S -wave states for the quarkonium with the Martin potential by the Bohr-Sommerfeld procedure.

Moreover, with the accuracy up to the logarithmic corrections, α_S is the constant value. Thus, as it has been shown in refs.[4, 5], for the leptonic constants of S -wave quarkonia, the scaling relation takes place

$$\frac{f^2}{M} \left(\frac{M}{4\mu} \right)^2 = \text{const.} , \quad (47)$$

independently of the heavy quark flavours.

Taking into the account eqs.(43) and (44) and integrating eqs.(39), (40) by parts, one can get with the accuracy up to border terms, that one has

$$-2f_n \frac{df_n}{dn} \frac{dn}{dM_n} n = \frac{\alpha_s}{\pi} M_n \left(\frac{4\mu}{M_n} \right)^2 . \quad (48)$$

Comparing eqs.(42) and (48), one finds

$$\frac{df_n}{f_n dn} = -\frac{1}{2n} , \quad (49)$$

that gives, after the integration, eq.(3):

$$\frac{f_{n_1}^2}{f_{n_2}^2} = \frac{n_2}{n_1} .$$

Table 1: The experimental values of leptonic constants (in MeV) for the nS -bottomonia in comparison with the estimates of present model.

quantity	exp.	present
f_1	715 ± 15	input
f_2	487 ± 16	506 ± 10
f_3	429 ± 14	412 ± 8
f_4	320 ± 30	358 ± 7
f_5	369 ± 46	320 ± 7
f_6	240 ± 30	292 ± 6

Relation (3) leads to that the border terms, which have been neglected in the writing of eq.(48), are identically equal to zero.

Thus, under the conditions of small Λ_{QCD}/m_Q ratio value and the nonrelativistic heavy quark motion, the described QCD sum rule scheme with the "smooth average value" allows one to derive the scaling expression, relating the leptonic constants of different nS -wave levels, independently of the quark content of heavy quarkonium.

2 Analysis of Scaling Relation

First, note that eq.(47), relating the leptonic constants of different quarkonia, turns out to be certainly valid for the quarkonia with the hidden flavour ($c\bar{c}$, $b\bar{b}$), where $4\mu/M = 1$ (see [4]). In that case, to estimate the constant value in the right hand side of eq.(47) we have supposed in refs.[4, 5], that α_S has the value, defining the Coulomb part of potential in the Cornell model (see [2]), and the $\langle dM/dn \rangle$ value is equal to the average distance between the nearest S -wave levels in the bottomonium.

Second, eq.(47) gives estimates of the leptonic constants for the heavy B and D mesons, so these estimates are in a good agreement with the values, obtained in the framework of other schemes of the QCD sum rules [3].

These two facts show that the offered scheme can be quite reliably applied to the systems with the heavy quarks.

Taking a value of the $1S$ -level leptonic constant as the input one, we

calculate the leptonic constants of higher nS -excitations in the charmonium and the bottomonium.

The results are presented in Tables 1, 2 and on Figures 1, 2. One can see that eq.(3) is in a good agreement with the experimental values of leptonic constants for the nS -wave levels of heavy quarkonia [9].

One has to note, that for the $f_{\psi(3S)}$ value we have taken

$$f_{\psi(3S)}^2 = f_{\psi(3770)}^2 + f_{\psi(4040)}^2 , \quad (50)$$

since, as it generally accepted, the $\psi(3770)$ and $\psi(4040)$ states are the results of $3D$ - and $3S$ -levels mixing in the charmonium.

Conclusion

In the framework of the QCD sum rules for the leptonic constants of the heavy quarkonia one uses the conditions of low Λ_{QCD}/m_Q ratio value and the non-relativistic quark motion in the phenomenological potential, possessing the simple scaling properties, one takes into the account the Coulomb-like α_S/v corrections, and in the scheme of the "smooth average value" one derives the scaling expression, relating the leptonic constants of nS -wave quarkonium levels, so

$$\frac{f_{n_1}^2}{f_{n_2}^2} = \frac{n_2}{n_1} ,$$

independently of the heavy quark flavours.

Table 2: The experimental values of leptonic constants (in MeV) for the nS -charmonia in comparison with the estimates of present model.

quantity	exp.	present
f_1	410 ± 14	input
f_2	283 ± 14	290 ± 10
f_3	205 ± 20	237 ± 8
f_4	180 ± 30	205 ± 7
f_5	145 ± 15	183 ± 6

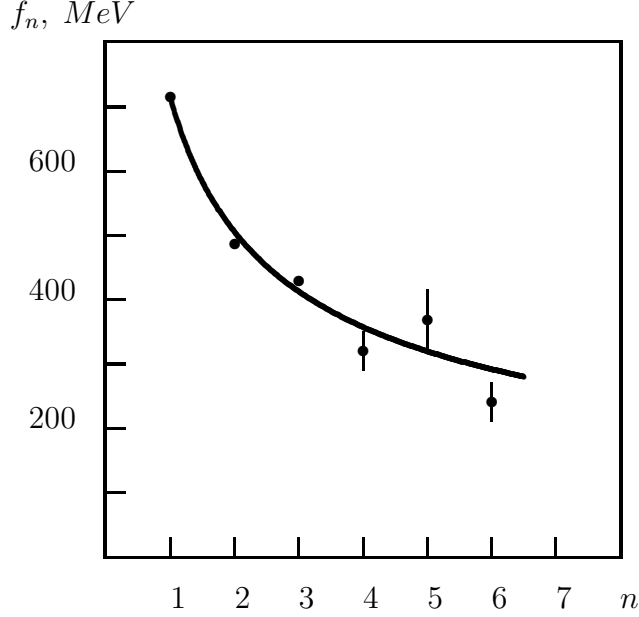


Figure 1: The calculated dependence of nS - bottomonium leptonic constants and the experimental values of $f_{\Upsilon(nS)}$.

The obtained relation is in a good agreement with the experimental values of leptonic constants for the charmonium and the bottomonium, and it reflects a small variation of the heavy quark kinetic energy with respect to the heavy quark flavours.

References

- [1] S.Nussinov and W.Wentzel, Phys.Rev. D36 (1987) 130;
M.B.Voloshin and M.A.Shifman, Sov.J.Nucl.Phys. 45 (1987) 292, 47 (1988) 511;
G.P.Lepage and B.A.Thacker, Nucl.Phys. B4 (proc.Suppl.) (1988) 199;
E.Eichten, Nucl.Phys. B4 (Proc.Suppl.) (1988) 170;
H.D.Politzer and M.B.Wise, Phys.Lett. B206 (1988) 681, B208 (1988) 504;

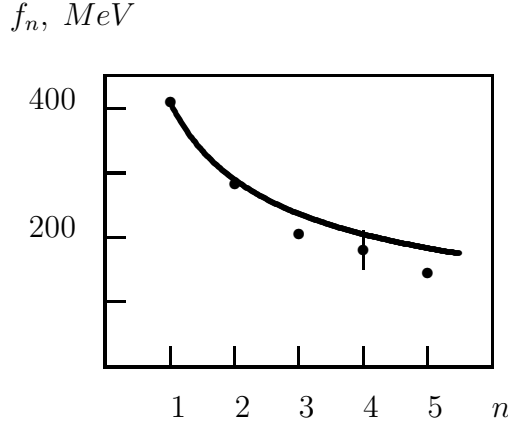


Figure 2: The calculated dependence of nS - charmonium leptonic constants and the experimental values of $f_{\psi(nS)}$.

N.Isgur and M.B.Wise, Phys.Lett. B232 (1989) 113, B237 (1990) 527, Phys.Rev.Lett. 66 (1991) 1130, Nucl.Phys. B348 (1991) 278, Phys.Rev. D4 (1990) 2388, D43 (1991) 819.

E.Eichten and B.Hill, Phys.Lett. B234 (1990) 511;

H.Georgi, Phys.Lett. B240 (1990) 447;

B.Grinstein, Nucl.Phys. B339 (1990) 253;

J.D.Bjorken, SLAC-PUB-5278, invited talk at Rencontre de Physique de la Vallée d'Acoste, La Thuile, Italy (1990);

M.Suzuki, Nucl.Phys. B258 (1985) 553;

B.Grinstein et al., Phys.Rev.Lett. 56 (1986) 298;

T.Altomari and L.Wolfenstein, Phys.Rev.Lett. 58 (1987) 1583;

N.Isgur, D.Scora, B.Grinstein and M.B.Wise, Phys.Rev. D39 (1989) 799;

E.V.Shuryak, Phys.Lett. 93B (1980) 134, Nucl.Phys. B198 (1982) 83, B32 (1989) 799;

A.Falk et al., Nucl.Phys. B343 (1990) 1.

[2] E.Eichten and F.Feinberg, Phys.Rev.Lett. 43 (1979) 1205, Phys.Rev. D23 (1981) 2724;

E.Eichten et al., Phys.Rev. D21 (1980) 203;

J.L.Richardson, Phys.Lett. 82B (1979) 272;

W.Buchmüller and S.-H. H.Tye, Phys.Rev. D24 (1981) 132.

- [3] M.A.Shifman, A.I.Vainstein, V.I.Zakharov, Nucl.Phys. B147 (1979) 385, 448;
 L.J.Reinders, H.Rubinstein, T.Yazaki, Phys.Rep. 127 (1985) 1;
 V.A.Novikov et al., Nucl.Phys. B237 (1989) 525;
 E.V.Shuryak, Nucl.Phys. B198 (1982) 83;
 T.M.Aliev, V.L.Eletsy, Yad.Fiz. 38 (1983) 1537 [Sov.J.Nucl.Phys. 38 (1983) 936];
 M.A.Shifman, Usp.Fiz.Nauk 151 (1987) 193 [Sov.Phys.Uspeski 30 (1987) 91];
 L.J.Reinders, Phys.Rev. D38 (1988) 947;
 S.Narison, Phys.Lett. B198 (1987) 104;
 C.A.Dominguez and N.Paver, Phys.Lett. B197 (1987) 423; B199 (1987) 596.
- [4] V.V.Kiselev, Nucl.Phys. B406 (1993) 340.
- [5] V.V.Kiselev, Preprint IHEP 94-63, Protvino, 1994.
- [6] E.Eichten, Preprint FERMILAB-Conf-85/29-T, 1985.
- [7] C.Quigg and J.L.Rosner, Phys.Lett. B71 (1977) 153.
- [8] A.Martin, Phys.Lett. 93B (1980) 338.
- [9] K.Hikasa et al., Particle Data Group, Phys.Rev. D45 (II) (1992) S1.